CERN-TH/95-261

Physics at LEP and Yukawa Coupling Unification*

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Abstract

We discuss the impact of the recent precision measurements at LEP on the minimal supersymmetric standard model. We show that the values of $\tan\beta$ necessary to induce large positive corrections to R_b are the same as the ones preferred by the unification of the bottom and τ Yukawa couplings in the MSSM for the current measured value of M_t. We discuss the relevance of the preferred parameter space for the Higgs and sparticle searches at LEP2. Remarkably, it follows that the LEP measurements can provide information about the structure of soft supersymmetry breaking parameters at M_{GUT}. Finally, we briefly discuss the properties of a supersymmetric model with four generations, for which the fit to R_b naturally improves with respect to the one in the Standard Model.

CERN-TH/95-261
October 1995

* Talk presented at **SUSY95**, Palaiseau, France, May 1995. To appear in the Proceedings.

One of the strongest motivations for the analysis of grand unified scenarios, which lead to the minimal supersymmetric standard model (MSSM) as an effective low-energy theory, comes from the successful prediction for the weak mixing angle within this framework [1]. Apart from yielding relations between the gauge couplings, most grand unified scenarios relate the values of the third-generation Yukawa couplings at M_{GUT}, leading to definite low-energy predictions for these quantities. The recent measurements of the top quark mass [2] at the Tevatron collider allow a restriction of the parameter space preferred by bottom-τ Yukawa coupling unification to two distinct regions: the region of very low $\tan\beta$, associated with the infrared fixed point for the top quark mass, and the region of very large $\tan\beta \simeq m_t/m_b$ [4, 5, 6]. It is the purpose of this talk, in the light of the most recent precision measurements at LEP, to investigate the most relevant experimental signatures associated with the theoretical scenario of Yukawa coupling unification.

The electroweak precision measurements at LEP and SLD are in remarkable agreement with the predictions of the Standard Model, in the presence of a heavy top quark, $M_t \simeq 175$ GeV [3]. The most recent experimental data show, however, that the partial quark width R_b (R_c) [$R_q = \Gamma(Z \rightarrow q\bar{q})/\Gamma(Z \rightarrow \text{hadrons})$], lies more than three (two) standard deviations above (below) the standard model fit for these quantities. If these measurements were taken at face value, they would point towards the presence of new physics at the weak scale, which can lead to large $Zq\bar{q}$ vertex corrections. From the experimental point of view, however, the measurement of R_c is subject to large experimental uncertainties. Moreover, if new physics were responsible for the deviation in R_c it should affect also the total light quark widths in a very precise way not to spoil the predictions for the total hadron width. The measurement of R_b is instead more reliable; hence, the present discrepancy with the standard model prediction deserves further investigation. The LEP electroweak working group provides also a determination of R_b based on the assumption that R_c is approximately given by its standard model value: $R_b = 0.2205 \pm 0.0016$, a value which is still more than three σ above the standard model prediction for this quantity: $R_b^{SM} \simeq 0.2156$.

Within the minimal supersymmetric standard model, large corrections to R_b are always associated with regions of large top and/or bottom Yukawa couplings [7, 8]. These couplings appear in the chargino-stop (charged Higgs-top) and neutralino-sbottom (neutral Higgs-bottom) one-loop contributions to the $Zb\bar{b}$ vertex. The value of the top Yukawa coupling is maximized for low values of $\tan\beta$, and hence large corrections to R_b are observed close to the top quark mass infrared fixed-point solution [9]. Since the charged Higgs contributions reduce the value of R_b, large positive corrections may only be obtained for large values of the charged Higgs mass. Stop-chargino loops give, instead, positive contributions to R_b; thus low values for the stop mass and, most important, for the chargino mass, within the reach of LEP2, are preferred. Large bottom Yukawa couplings, instead, may only be obtained for large values of $\tan\beta \simeq m_t/m_b$. In this region of parameters, the total Higgs contributions become positive for $m_A \lesssim 70$ GeV, and for sufficiently low values of m_A they are larger than the genuine supersymmetric ones [8, 10]. Low values of the third-generation squark, chargino and neutralino masses are also helpful in order to get large radiative corrections to R_b.

It is remarkable that the regions of $\tan\beta$ preferred by the best fit to precision measurement data and by the requirement of bottom-τ Yukawa unification coincide. Let us first concentrate in the low $\tan\beta$ region, associated with the infrared fixed point of the top quark mass. One of the interesting properties of the top quark mass infrared fixed-point solution is that it implies also a fixed-point solution for some soft supersymmetry-breaking parameters [11]. Summarizing the results for the relevant low-energy mass parameters at the infrared fixed point we have [9]:

$$
\begin{aligned}
m_{H_2}^2 &\simeq \frac{m_{H_2}^2(0)}{2} - m_Q^2(0) - 3.5 M_{1/2}^2 \qquad\qquad A_t \simeq -2M_{1/2} \\
m_Q^2 &\simeq \frac{2m_Q^2(0)}{3} - \frac{m_{H_2}^2(0)}{6} + 6M_{1/2}^2 \qquad\qquad m_U^2 \simeq \frac{m_Q^2(0)}{3} - \frac{m_{H_2}^2(0)}{3} + 4M_{1/2}^2 \\
\mu^2 &\simeq \left[m_{H_1}^2(0) + \left(\frac{2m_Q^2(0) - m_{H_2}^2(0)}{2} \right) \tan^2\beta + M_{1/2}^2 \left(0.5 + 3.5 \tan^2\beta \right) \right] \frac{1}{\tan^2\beta - 1}
\end{aligned}
\tag{1}
$$

where $m_i^2(0)$ and $M_{1/2}$ denote the values of the scalar mass parameters and the gaugino masses at M_{GUT} (for notation and sign conventions, see Ref. [9]). Observe that A_t and $M_U^2 = m_Q^2 + m_U^2 + m_{H_2}^2$ are determined by the parameters of the gauge sector of the theory.

Another very important feature of the spectrum at the infrared fixed point is associated with the Higgs sector. The Higgs spectrum consists of three neutral scalar states –two CP-even, h and H, and one CP-odd, A– and two charged scalar states H$^{\pm}$. The mass of the CP-odd scalar Higgs is approximately given by its tree level expression, $m_A^2 \simeq m_1^2 + m_2^2$,

$$
m_A^2 \simeq \left[m_{H_1}^2(0) + \left(\frac{2m_Q^2(0) - m_{H_2}^2(0)}{2} \right) + 4M_{1/2}^2 \right] \frac{\tan^2\beta + 1}{\tan^2\beta - 1} \tag{2}
$$

From Eq. (2) it follows that, for lower values of $\tan\beta$, the value of the CP-odd mass eigenstate is enhanced. Moreover, larger values of m_A imply as well that the charged Higgs and the heaviest CP-even Higgs become heavy in such a regime. For these large values of m_A, m_h acquires values close to its upper bound, which is independent of the exact value of the CP-odd mass. Furthermore, for a given value of the physical top quark mass, the infrared fixed-point solution is associated with the minimum value of $\tan\beta$ compatible with the perturbative consistency of the theory. For $\tan\beta \geq 1$, lower values of $\tan\beta$ correspond to lower values of the tree-level lightest CP-even mass, $m_h^{tree} = M_Z|\cos 2\beta|$, and after the inclusion of the radiative corrections, the fixed point still gives the lowest possible value of m_h for any given value of M_t [12, 13]. Indeed, taking into account the most relevant two-loop corrections [14], an approximate expression for the Higgs mass may be obtained [15],

$$
\begin{aligned}
m_h^2 &= M_Z^2 \cos^2 2\beta \left(1 - \frac{3}{8\pi^2} \frac{m_t^2}{v^2}\, t \right) \\
&+ \frac{3}{4\pi^2} \frac{m_t^4}{v^2} \left[\frac{1}{2} \tilde{X}_t + t + \frac{1}{16\pi^2} \left(\frac{3}{2} \frac{m_t^2}{v^2} - 32\pi\alpha_3(M_t) \right) \left(\tilde{X}_t t + t^2 \right) \right]
\end{aligned}
\tag{3}
$$

where $\tilde{X}_t = 2\tilde{A}_t^2/M_{\rm SUSY}^2 \left(1 - \tilde{A}_t^2/12M_{\rm SUSY}^2 \right)$, $\tilde{A}_t = A_t - \mu\cot\beta$, $t = \log\frac{M_{\rm SUSY}^2}{M_t^2}$, $m_t = m_t(M_t)$ and $v = 174$ GeV. Taking the value of $\tilde{X}_t$ which maximizes the Higgs mass, and for $M_{SUSY} \lesssim 1$ TeV and $M_t \lesssim 175$ GeV, an upper bound $m_h \lesssim 100$ GeV is obtained. The most recent experimental analyses have shown that for such range a of values of M_t the infrared fixed point solution can be fully tested at LEP2 [16].

A question that immediately arises is that of the maximal value of R_b, which may be obtained for arbitrary choices of the boundary conditions for the mass parameters. The larger variations of R_b with respect to its Standard Model value are found for solutions such that the right-handed component of the lightest stop is maximized by requiring low (large) values for the mass parameter m_U (m_Q). These conditions also imply that the stop contributions to the variable $\Delta\rho$ are small. Moreover, the lightest stop and chargino masses and the μ parameter should acquire values below M_Z. Finally, the charged Higgs must be heavy, $m_{H^+} \gg M_Z$. Light stops demand significant mixing effects, $A_t \simeq m_Q$. Taking $M_t \simeq 160$–180 GeV ($\tan\beta \simeq 1.1$–1.6) and imposing the experimental constraints coming from $\Delta\rho$ and the branching ratio $b \to s\gamma$, the maximum value of R_b achievable is $R_b \simeq 0.2180$. This value of R_b has the effect of changing the strong gauge coupling determination at LEP, coming from $R_\ell = \Gamma_\ell/\Gamma_h$, where Γ_ℓ is the Z leptonic width, to a value $\alpha_3(M_Z) \simeq 0.116$.

Having found these solutions, it is important to analyse for which particular values of the boundary conditions for the scalar mass parameters they are obtained. The set of values [9]

$$m_Q^2(0) = m_U^2(0) \simeq 0; \qquad m_{H_2}^2(0) \simeq 12 M_{1/2}^2; \qquad m_{H_1}^2(0) \simeq 2.5\ \tan^2\beta\ M_{1/2}^2, \tag{4}$$

with $M_{1/2} > 300$ GeV, leads to solutions in the desired range. The first two relations in Eq. (4) are necessary in order to minimize the value of m_U while keeping $A_t \simeq m_Q$. The last relation is required in order to minimize the value of the μ parameter. Interestingly, the solutions that maximize R_b demand a small value of the squark mass parameters at the grand unification scale. Finally, the charged Higgs mass is naturally large within this scheme.

We shall now discuss the large $\tan\beta$ region [17], for which, as explained above, low values of the CP-odd Higgs mass, $m_A \lesssim 70$ GeV, are preferred to improve the agreement with the value of R_b measured at LEP. An important tree level relation is obtained for large values of the ratio of Higgs vacuum expectation values, $\tan\beta > 10$, for which the tree level value of the lightest CP-even Higgs mass is equal to M_Z, whenever the CP-odd Higgs mass m_A is larger than M_Z, while for $m_A \leq M_Z$, $m_h = m_A$. This tree level relation is approximately stable under radiative corrections, with the only difference that the range for which the equality $m_h = m_A$ holds, extends to values of m_A somewhat larger than M_Z. Therefore, large values of $\tan\beta$ and values of the CP-odd Higgs in the desired range imply

$$m_h < M_Z, \tag{5}$$

and hence the neutral Higgs sector will be tested at LEP2 within this framework.

Moreover, the charged Higgs mass is determined through the CP-odd Higgs mass value, $m_{H^\pm}^2 \simeq m_A^2 + M_W^2$. For $m_{H^\pm} \lesssim 130$ GeV, the branching ratio BR($b \to s\gamma$) becomes larger than the presently allowed experimental values, unless the supersymmetric particle contributions suppress the charged Higgs enhancement of the decay rate. The most important supersymmetric contributions to this rare bottom decay come from the chargino-stop one-loop diagram [24]. The chargino contribution to the $b \to s\gamma$ decay amplitude depends on the soft supersymmetry breaking mass parameter A_t and on the supersymmetric mass parameter μ, for very large values of $\tan\beta$, it is given by

$$A_{\tilde{\chi}^+} \simeq \frac{m_t^2}{m_{\tilde{t}}^2} \frac{A_t \mu}{m_{\tilde{t}}^2} \tan\beta \, G\left(\frac{m_{\tilde{t}}^2}{\mu^2}\right), \tag{6}$$

where $G(x)$ is a function that takes values of order 1 when the characteristic stop mass $m_{\tilde{t}}$ is of order μ and grows for low values of μ. One can show that for positive (negative) values of $A_t \times \mu$ the chargino contributions are of the same (opposite) sign as the charged Higgs ones. Hence, to get an acceptable $b \to s\gamma$ decay rate, negative values for $A_t \times \mu$ are required. Considering the running of the soft supersymmetry-breaking terms from high energies, the value of A_t is approximately given by $A_t = A_0 \left(1 - \frac{Y_t}{Y_f}\right) - M_{1/2}\left(4 - 2\frac{Y_t}{Y_f}\right)$ [18], where A_0 and $M_{1/2}$ are the values of the trilinear coupling and the common gaugino mass at the unification scale and Y_f is the fixed-point value for Y_t. For values of the top quark mass $M_t \gtrsim 160$ GeV, the mass parameter A_t is opposite in sign to $M_{1/2}$, unless A_0 is more than one order of magnitude larger than $M_{1/2}$. Hence, in order to get acceptable values for the branching ratio BR($b \to s\gamma$),

$$\mu \times M_{1/2} > 0. \tag{7}$$

Large radiative corrections to the bottom mass arise from the effective coupling of the bottom quark to the Higgs H_2 at the one loop level [19, 18]. Indeed, although the effective coupling of the down quarks to H_2 is small in comparison to the down quark Yukawa coupling to the Higgs H_1, for large values of $\tan\beta$ the vacuum expectation value of H_2 is much larger than that of H_1, and hence large down quark mass corrections may be obtained through this effect [20]. Quantitatively, $\Delta m_b/m_b \simeq \tan\beta \times 2\alpha_3 M_{\tilde{g}}\mu/(3\pi m_{\tilde{b}}^2)$, where $Y_t = h_t^2/(4\pi)$ and $m_{\tilde{b}}^2$ is the value of the heaviest sbottom mass (for more precise expressions see Refs. [19, 18, 20]). Due to the constraints on the parameters μ and $M_{1/2}$ discussed above, the bottom mass corrections are positive within this scheme. Positive mass corrections mean that, in order to get an acceptable values for the physical bottom mass, the value of the predicted bottom mass before addition of the finite corrections, $\tilde{M}_b$, should be given by $\tilde{M}_b \leq 5.2$GeV. This puts strong constraints on the allowed values of M_t. For $\alpha_3(M_Z) \gtrsim 0.125, 0.120, 0.115$, the top quark mass $M_t \gtrsim 180, 170, 160$ GeV (See Ref. [18]).

Finally, let us discuss the model with four generations [25]: If one considers the R_b measurement alone within the context of the standard model, one would conclude that $m_t \lesssim m_W$. Previous top quark searches at hadron colliders are able to close the window between the top mass lower bound coming from the $W^\pm$ width, 62 GeV, and $m_W + m_b \simeq 85$ GeV, assuming that $t \to bW^\star$ is the dominant top-quark decay mode. However, if the top quark had any two-body decay modes (due to new physics processes), and if these modes rarely produced leptons, then a top quark in this mass region would not have been detected in any experiment. A well motivated and experimentally acceptable scenario occurs in supersymmetric models in which the decay $t \to \tilde{t}\tilde{\chi}_1^0$ is kinematically allowed (where $\widetilde{t}$ is the lightest top squark and $\widetilde{\chi}_1^0$ is the lightest neutralino [26]). We choose $m_t \simeq m_W$, $M_{\tilde{t}} \simeq 50$ GeV and $M_{\tilde{\chi}_1^0} \simeq 25$ GeV. Then, the dominant decay chain is $t \to \widetilde{t}\widetilde{\chi}_1^0$ followed by $\widetilde{t} \to c\widetilde{\chi}_1^0$ through a one-loop process, which rarely produces a hard isolated lepton.

The fourth-generation top quark decays $t' \to bW^+$ can be the source of the CDF and D0 events. For this, the t'–b mixing angle ($V_{t'b}$) must not be too small; otherwise, the $t' \to b' + W^+$ decay mode will dominate. For definiteness, we choose $|V_{t'b}/V_{t'b'}| = 0.1$. It is straightforward to show that, for $m_{t'} \simeq 175$ GeV, a BR($t' \to bW^+$) close to 1 is obtained if $m_{b'} \gtrsim 105$ GeV. Moreover, imposing the requirement that the Yukawa couplings remain perturbative up to the grand unification scale, we obtain that two properties must be fulfilled: first, the b' mass must be close to its lower bound, $m_{b'} \simeq 105$ GeV. Second, the b' and t' Yukawa couplings must lie close to their fixed-point values and $\tan\beta \simeq 1.6$.

The fourth-generation lepton Yukawa couplings are also strongly constrained by the experimental bound on the lepton masses. The bound on the τ' lepton mass implies that its Yukawa coupling must get large values at the grand unification scale. Assuming that the fourth-generation neutrino (N) is a Dirac fermion, one can also add the requirement that its Yukawa coupling takes large values at M_{GUT}. The resulting lepton masses are: $m_{\tau'} \simeq 50$ GeV and $m_N \simeq 80$ GeV. Remarkably, the fourth generation fermion masses within this model are consistent with the unification of the four fourth generation Yukawa couplings at the GUT scale.

The constraints coming from the oblique corrections to the gauge bosons, most notably the $\Delta\rho$ parameter, put strong constraints on the sparticle masses. Indeed, since the top quark mass is less than half of its standard value, the contribution of the t–b doublet to $\Delta\rho$ is reduced by a factor of 4. Due to the large values of $m_{b'}$, the fourth-generation quark contribution to $\Delta\rho$ is approximately the same as the third-generation one. Hence, the supersymmetric particle virtual effects must account for roughly half of the total contribution to $\Delta\rho$. One must maximize the off-diagonal squark mixing while keeping the diagonal squark mass parameters as small as possible. However, the latter cannot be too small; otherwise the radiative corrections to the light Higgs mass will be reduced leading to a value of m_h below the current LEP bound.

For $m_Q \simeq m_U$, relevant contributions to $\Delta\rho$ induced by the third-generation squarks may only be obtained if m_Q is somewhat above the lower bounds on this quantity coming from the present experimental bounds on the lightest sbottom mass, implying large values of the stop mixing mass parameters. Non-negligible mixing in the fourth-generation also enhances the fourth generation squark contributions to $\Delta\rho$. The maximum effect is limited phenomenologically by a lower bound on the mass of $\widetilde{b}'$. In order that $t' \to bW^+$ remain the dominant decay, one must kinematically forbid $t' \to \widetilde{b}'\widetilde{\chi}_1^+$. This in turn determines the best value for μ, since $A_{t'}$ and $A_{b'}$ are determined by their infrared fixed-point behaviour [11, 18], $A_{t'} \simeq A_{b'} \simeq -200$ GeV. Taking this into account we find that, for instance, if $m_Q \simeq m_U \simeq 275$ GeV, $\mu \simeq -420$ GeV $m_{Q'} \simeq m_{U'} \simeq 250$ GeV, and $A_t \simeq 750$ GeV, the contribution from the third- and fourth- generation squarks is only slightly lower than that one of their fermion superpartners, leading to acceptable values for $\Delta\rho$. The price to pay is a very large value of the stop mixing mass term $\tilde{A}_t \simeq 1$ TeV. These large values of the mixing mass parameter $\tilde{A}_t$ may render the ordinary vacuum state unstable [9, 21, 22]. Our model predicts a light Higgs spectrum, $m_h \simeq 70$ GeV, and an improved value of $R_b = 0.2184$, which is within 1.5σ of the measured LEP value.

In summary, the three scenarios discussed above are consistent with Yukawa coupling unification. In each scenario the global fit to the present electroweak precision measurement is significantly improved if new particles, with masses below the LEP2 kinematical limit, are present:

a) Infrared fixed point: A chargino, two neutralinos and a stop. Furthermore, the lightest CP-even Higgs (if $M_t \lesssim 175$ GeV), must be observable at LEP2.

b) Large $\tan\beta$: The CP-odd and the lightest CP-even Higgs bosons, a chargino, two neutralinos and a stop.

c) Model with four generations: The lightest CP-even Higgs, a charged and a neutral lepton, a chargino, two neutralinos, a stop and the top quark.

Acknowledgement: I would like to thank M. Carena, P. Chankowski, H. Haber, M. Olechowski, S. Pokorski and M. Quirós for fruitful and enjoyable collaborations.

References

[1] S. Dimopoulos, S. Raby and F. Wilczek, *Phys. Rev.* **D24** (1981) 1681;
S. Dimopoulos and H. Georgi, *Nucl. Phys.* **B193** (1981) 150;
L. Ibañez and G.G. Ross, *Phys. Lett.* **B105** (1981) 150.

[2] F. Abe et al., CDF Collaboration, *Phys. Rev.* **D50** (1994) 2966; *Phys. Rev. Lett.* **73** (1994) 225; preprint FERMILAB-PUB-95/022-E (2 March 1995)
S. Abachi et al., D0 Collaboration, *Phys. Rev. Lett.* **72** (1994) 2138; *Phys. Rev. Lett.* **74** (1995) 2422; preprint FERMILAB-PUB-95/028-E (3 March 1995)

[3] P. Antilogus *et al.* [LEP Electroweak Working Group], LEPEWWG/95-02 (1995).

[4] H. Arason, D. J. Castaño, B. Keszthelyi, S. Mikaelian, E. J. Piard, P. Ramond and B. D. Wright, Phys. Rev. Lett. 67 (1991) 2933;
S. Dimopoulos, L. Hall and S. Raby, Phys. Rev. Lett. 68 (1992) 1984, Phys. Rev. D45 (1992) 4192.

[5] V. Barger, M.S. Berger and P. Ohmann, *Phys. Rev. Lett.* **49** (1994) 4908;
P. Langacker and N. Polonsky, *Phys. Rev.* **D47** (1993) 4028; **D49** (1994) 1454;
M. Carena, S. Pokorski and C.E.M. Wagner, *Nucl. Phys.* **B406** (1993) 59;
W.A. Bardeen, M. Carena, S. Pokorski and C.E.M. Wagner, *Phys. Lett.* **B320** (1994) 110.

[6] M. Olechowski, S. Pokorski, *Phys. Lett.* **B214** (1988) 393;
B. Anantharayan, G. Lazarides and Q. Shafi, *Phys. Rev.* **D44** (1991) 1613;

[7] M. Boulware and D. Finnel, *Phys. Rev.* **D44** (1991) 2054.

[8] J.D. Wells, C. Kolda and G.L. Kane, *Phys. Lett.* **B338** (1994) 219;
P. Chankowski and S. Pokorski, Warsaw University preprint, IFT-95/5, to appear in the Proceedings of "Beyond the Standard Model IV", Lake Tahoe, CA, December 1994; MPI preprint MPI-PhT/95-49 (1995);
A. Dabelstein, W. Hollik and W. Mösle, Univ. of Karlsruhe preprint, KA-THEP-5-1995 (1995).

[9] M. Carena and C.E.M. Wagner, Nucl. Phys. **B452** (1995) 45.

[10] D. Garcia and J. Sola, *Phys. Lett.* **B354** (1995) 335.

[11] M. Carena, M. Olechowski, S. Pokorski and C.E.M. Wagner, *Nucl. Phys.* **B419** (1994) 213.

[12] V. Barger, M.S. Berger and P. Ohmann, *Phys. Rev.* **D47** (1993) 1093; V. Barger, M.S. Berger, P. Ohmann and R.J.N. Phillips, *Phys. Lett.* **B314** (1993) 351.

[13] C.E.M. Wagner, *Properties of SUSY particles*, eds. L. Cifarelli and V.A. Khoze (World Scientific, Singapore, 1993) p. 469.

[14] R. Hempfling and A.H. Hoang, *Phys. Lett.* **B331** (1994) 99;
J. Kodaira, Y. Yasui and K. Sasaki, *Phys. Rev.* **D50** (1994) 7035;
J.A. Casas, J.R. Espinosa, M. Quirós and A. Riotto, *Nucl. Phys.* **B436** (1995) 3; (E) **B439** (1995) 466.

[15] M. Carena, J.R. Espinosa, M. Quiros and C.E.M. Wagner, *Phys. Lett.* **B355** (1995) 209;
M. Carena, M. Quiros and C.E.M. Wagner, CERN preprint CERN-TH/95-157, August 1995, Submitted to *Nucl. Phys. B.*

[16] G. Altarelli et al., *Interim Report on the Physics Motivations for an Energy Upgrade of LEP2*, CERN preprint CERN-TH/95-151, CERN-PPE/95-78, June 1995.

[17] For a more extensive discussion, see M. Carena, P. Chankowski, M. Olechowski, S. Pokorski and C.E.M. Wagner, to appear.

[18] M. Carena, M. Olechowski, S. Pokorski and C.E.M. Wagner, *Nucl. Phys.* **B426** (1994) 269.

[19] L.J. Hall, R. Rattazzi and U. Sarid, *Phys. Rev.* **D50** (1994) 7048;
R. Hempfling, *Phys. Rev.* **D49** (1994) 6168.

[20] M. Carena, S. Dimopoulos, S. Raby and C.E.M. Wagner, CERN preprint CERN-TH/95-53, hep-ph/9503488, to be published in Phys. Rev. D;
T. Blažek, S. Raby and S. Pokorski, Ohio Univ. preprint OHSTPY-HEP-T-95-007, hep-ph/9504364, to be published in Phys. Rev. D.

[21] J.M. Frère, D.R.T. Jones and S. Raby, *Nucl. Phys.* **B222** (1983) 11;
L. Alvarez Gaumé, J. Polchinski and M. Wise, *Nucl. Phys.* **B221** (1983) 495;
C. Kounnas, A. Lahanas, D.V. Nanopoulos and M. Quirós, *Nucl. Phys.* **B236** (1984) 438;
J.F. Gunion, H.E. Haber and M. Sher, *Nucl. Phys.* **B306** (1988) 1.

[22] P. Langacker and N. Polonsky, *Phys. Rev.* **D50** (1994) 2199;
J.A. Casas, A. Lleyda and C. Muñoz, preprint FTUAM 95/11 (1995).

[23] B. Barish et al., CLEO Collaboration, preprint CLEO CONF 94-1, to appear in the Proceedings of the ICHEP94 Conference, Glasgow, Scotland, July 1994.

[24] S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, *Nucl. Phys.* **B353** (1991) 591;
R. Barbieri and G. Giudice, *Phys. Lett.* **B309** (1993) 86.

[25] For a more extensive discussion, see M. Carena, H.E. Haber and C.E.M. Wagner, CERN preprint CERN-TH/95-235, October 1995.

[26] I.I. Bigi and S. Rudaz, *Phys. Lett.* **B153** (1985) 335.